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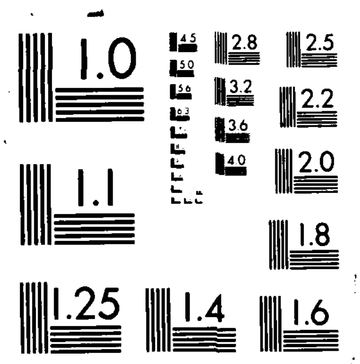
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Transverse Kelvin-Helmholtz Instability with Parallel Electron Dynamics and Coulomb Collisions

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Electron parallel dynamics and Coulomb collisions are included in the analysis of the transverse Kelvin-Helmholtz instability. The electrons are treated kinetically while the ions are treated in the fluid limit. It is shown that, in the collisionless case, for an inhomogeneous velocity profile $V(x) = V_0 \tanh(x/L)$ the Kelvin-Helmholtz instability is stable for $k_z/k_y > (V_0/L\,\omega_{\ell h})k_yL(1 - k_y^2L^2)^{1/2}$ in the limit $\omega - k_yV_0 >> k_zv_e$. Here, V_0 is the flow velocity, L is the scale length of the velocity shear layer, k_z and k_y are the parallel and perpendicular wavenumbers, respectively, $\omega_{\ell h} = (\Omega_e\Omega_i)^{1/2}$, and v_e is the electron thermal velocity. The stabilization of the mode is shown to be caused by the compressional energy given to the electrons parallel to $\underline{B}$. In the collisional limit, Coulomb collisions are shown to increase the unstable k_z domain because they inhibit the electron motion parallel to $\underline{B}$. Applications to the high latitude ionosphere are discussed.

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TRANSVERSE KELVIN-HELMHOLTZ INSTABILITY WITH PARALLEL ELECTRON DYNAMICS AND COULOMB COLLISIONS

I. INTRODUCTION

Velocity shear layers with steep gradient scale lengths are frequently observed in the high latitude auroral zone (Kelley and Carlson, 1982; Basu et al., 1985; Kelley and Earle, 1986). Recent HILAT satellite measurements reported by Basu et al. (1985) show flows with moderate to strong velocity shears with a shallow irregularity spectrum; often field aligned currents are also observed in the regions of sheared velocity flows. The observations suggest that the east-west $\underline{E} \times \underline{B}$ drift velocity, for example, is inhomogeneous and usually reverses its direction as one moves in the north-south direction. This kind of velocity shear transverse to the magnetic field is usually a source of Kelvin-Helmholtz instability (Chandrasekhar, 1961; Mikhailovskii, 1972). In the conventional Kelvin-Helmholtz instability the mixing of fluids with different velocities is confined to two dimensions transverse to the magnetic field; the ion polarization drift leads to the instability while the electrons simply $\underline{E} \times \underline{B}$ drift in the perpendicular plane. In the auroral ionosphere one cannot simply ignore the dynamics parallel to the magnetic field because there could be additional sources of free energy or damping that may affect the transverse mixing of the fluids with sheared velocity. Our aim is to introduce the third dimension into the two dimensional analysis of Kelvin-Helmholtz instability and, as a first step, we study the effects of parallel electron dynamics. Furthermore, since the electron parallel motion couples the collisionless and collisional regimes of the ionosphere, we also include Coulomb collisions between the electrons and ions. A preliminary analysis, in the collisionless domain, was performed by Thompson (1983) who showed that the parallel electron motion could have stabilizing influence on the Kelvin-Helmholtz instability. In this paper we perform (i) a more complete analysis and present a general

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stability criterion for a specific velocity profile, (ii) present the effects of collisions on the instability. In analogy with the role of gravity on Kelvin-Helmholtz instability in neutral fluids, we observe that the parallel electron dynamics simulate buoyancy of a fluid in a gravitational field.

II. THEORY

We consider a homogeneous plasma that is drifting across a magnetic field ($\underline{B} = B\,\hat{z}$) because of an equilibrium inhomogeneous electric field ($\underline{E} = E_0(x)\,\hat{x}$). We choose $E_0(x) = E_0 \tanh(x/L)$ where L is the scale length of the shear layer. In this paper we restrict the analysis to the domain $L \gg \rho_i$ and $|\omega| \ll \Omega_i$, where ρ_i and Ω_i are the ion gyroradius and gyrofrequency, respectively. Incompressible plasma motion in two dimensions across an uniform magnetic field is determined by the constraint that the perpendicular current be divergence free,

$$\nabla \cdot \underline{j}_\perp = 0 \tag{1}$$

where $\underline{j}_\perp = ne(\underline{v}_{e\perp} - \underline{v}_{i\perp})$, e is the charge, n is the plasma density, $\underline{v}_{e\perp}$ and $\underline{v}_{i\perp}$ are the perpendicular drifts of electrons and ions. The velocities $\underline{v}_{e\perp}$ and $\underline{v}_{i\perp}$ are obtained from the momentum equations by assuming the electrons only $\underline{E} \times \underline{B}$ drift, while the ions experience both the $\underline{E} \times \underline{B}$ drift and the polarization drift. Linearizing (1) with a perturbed electrostatic potential of the form $\hat{\Phi} \sim \hat{\phi}(x) \exp[-i(\omega t - k_y y)]$ with frequency ω and wavenumber k_y yields the equation governing the Kelvin-Helmholtz instabilty:

$$\frac{\partial^2 \hat{\phi}}{\partial x^2} + \left[- k_y^2 + \frac{k_y V_0''}{(\omega - k_y V_0)}\right]\hat{\phi} = 0 \tag{2}$$

where $V_0(x) = -(c/B)E_0(x)$. This is a well studied equation (Chandrasekhar, 1961; Mikhailovskii, 1964; Guzdar et al., 1982; Satyanarayana et al., 1983); instability ($\gamma > 0$, where γ is given by $\omega = \omega_r + i\gamma$) occurs in the wavenumber domain $0 < k_y L < 1$ for $V_0(x) = V_0 \tanh(x/L)$.

To extend the above theory to three dimensions, we need to consider the equation

$$\nabla_\perp \cdot \underline{j}_\perp = - \partial j_z/\partial z \tag{3}$$

which allows for a parallel current along the field lines. In this paper we consider perturbations along the magnetic field and also equilibrium currents along the field lines. We include electron collisions and use a BGK model to calculate the electron susceptibility (Clemmow and Dougherty, 1969). Kinetic treatment of the electron parallel motion yields a general equation for the perturbed electrostatic potential $\hat{\Phi} \sim \hat{\phi}(x) \exp[-i(\omega t - k_z z - k_y y)]$ in the domain $|\omega|$ and $V_0/L \ll \Omega_i$

$$\frac{\partial^2 \hat{\phi}}{\partial x^2} + \left[- k_y^2 + \frac{k_y V_0''}{(\omega + i\nu_{in} - k_y V_0)}\right]\hat{\phi} - \frac{\omega_{\ell h}^2}{v_e^2}\left[\frac{1 + \xi_e Z(\xi_e)}{1 + i(\nu_e/k_z v_e) Z(\xi_e)}\right]\hat{\phi} = 0 \tag{4}$$

where $\omega_{\ell h}^2 = \Omega_e \Omega_i$, $\Omega_e(\Omega_i)$ is the electron (ion) gyro-frequency, v_e is the electron thermal velocity, $\xi_e = (\omega - k_y V_0 - k_z V_d + i\nu_e)/k_z v_e$, , V_d is the equilibrium electron drift along the magnetic field, and ν_e is the electron collision frequency. For $\nu_e = 0$, (4) is the same as equation (8) of Thompson (1983) without the ion terms; we can ignore the parallel ion terms if we restrict our analysis to $\rho_i/L < 1$ where ρ_i is the ion gyroradius. We have introduced the parallel electron drift, V_d, and collisional effects which are absent in Thompson (1983).

III. COLLISIONLESS REGIME

(a) Analytical Results

In this section we set $\nu_e = 0$. A simple stability boundary defining the regions of k_z and k_y where the system is stable or unstable can be easily obtained in the domain where the electrons behave as a fluid; this is the region where the argument of the Z function is large (i.e., $|\xi_e| \gg 1$). In this domain the Z function can be approximated as $- 1/\xi_e - 1/2\ \xi_e^3$. With this approximation (4) takes the form

$$\frac{\partial^2 \hat{\phi}}{\partial x^2} + \left[- k_y^2 + \frac{k_y V_0''}{(\omega - k_y V_0)} + \frac{1}{2}\frac{k_z^2 \omega_{\ell h}^2}{(\omega - k_y V_0)^2}\right]\hat{\phi} = 0. \tag{5}$$

Equation (5) is identical to the Rayleigh equation governing the interchange of a drifting, weakly inhomogeneous fluid in a gravitational field with a sheared drift velocity (Drazin, 1958; Chandrasekhar, 1961). Here we identify the term $\omega_{\ell h}^2 k_z^2/k_y^2$ with the buoyancy term that is driving the Rayleigh-Taylor instability g/L_n, where g is gravity and L_n is the density gradient scale length. In the spirit of the neutral fluid problem, we define a dimensionless number

$$J = \frac{1}{2} \frac{k_z^2}{k_y^2} \frac{L^2\omega_{\ell h}^2}{V_0^2} . \tag{6}$$

Once we have made this identification, it follows that for $V_0(x) = V_0 \tanh(x/L)$ the stability boundary is simply (Drazin, 1958)

$$J = k_y^2L^2(1 - k_y^2L^2). \tag{7}$$

The modes are stable for $J > 1/4$ or

$$\frac{k_z^2}{k_y^2} > \frac{1}{2} \left(\frac{V_0}{L\omega_{\ell h}}\right)^2 . \tag{8}$$

Thus, parallel electron dynamics has a stabilizing influence on the Kelvin-Helmholtz instability.

Figure 1 shows the stability boundary where we plot J versus k_yL. We note that the conclusions of Thompson (1983) are qualitatively correct. However, our analysis shows that $J < 1/4$ for instability as opposed to $J < 1$ given by Thompson (1983) who did not consider the explicit shear profile $V(x) = V_0 \tanh(x/L)$. A derivation of this stability criterion based upon the energy principle is given in Sec. III.c.

(b) Numerical Results

We now solve (4) numerically to further show the effects of the parallel dynamics. In Figure 2 we plot the normalized growth rate $\gamma/(V_0L)$ vs. α $(= \omega_{\ell h}^2L^2/v_e^2)$ for $k_yL = 0.5$ (which gives maximum growth for $k_z = 0$), $k_z/k_y = 1.0 \times 10^{-3}$, $V_0/v_e = 0.01$, and $V_d = 0$. Curve A shows the growth rate for the collisionless case, while Curve B shows the growth rate when electron collisions are included (to be discussed in Sec. IV). For $\alpha = 0$ we find $\gamma/(V_0/L) = 0.19$ which is the usual result. As α is increased, we see the Kelvin-Helmholtz growth rate drops. Since for large α the Kelvin-Helmholtz term (i.e. $\propto V_0''$) does not contribute to (4), the mode equation essentially becomes

$$- k_y^2v_e^2 - \Omega_e\Omega_i[1 + \xi_eZ(\xi_e)] = 0 \tag{9}$$

which allows only stable roots in the fluid limit with $\omega_r - k_y V_0 = (1/2)(k_z/k_y)\omega_{\ell h}$. From Curve A of Figure 2 we also see that the modes are stable for $\alpha \gtrsim 9$. This is in agreement with the criterion $J > 1/4$ previously discussed. We find that for reasonable values of the parallel drift velocity (i.e., $V_d < v_e$) these results are not significantly altered. We find that in this regime the argument of the Z function is large, indicating that the electrons are merely acting as a fluid and wave particle effects are not playing an important role. On the other hand, for sufficiently small α the theory breaks down as shown by the kink in Fig. 2; the orbits begin to play an important role here ($L \sim \rho_i$) and kinetic effects need to included.

We now solve (4) for several mode numbers. Figure 3 shows the normalized growth rate $\gamma/(V_0 L)$ as a function of $k_y L$ for $V_0/v_e = 0.01$, $\alpha = 1$, and $k_z L = 0$, 1.0×10^{-3}, 2.0×10^{-3}, 3.0×10^{-3}, and 4.0×10^{-3} as curves A, B, C, D and E, respectively. Curve A shows the conventional Kelvin-Helmholtz mode; $\gamma > 0$ in the wavenumber domain $0 < k_y L < 1$ and $\gamma_{max} = 0.19$ for $k_y L = 0.46$. Several points are to be noted from this figure: (i) as $k_z L$ is increased the wavenumber of the maximally growing mode increases; for example, for $k_z L = 4.0 \times 10^{-3}$ the growth rate maximizes at $k_y L \sim 0.8$, (ii) the growth steadily decreases as $k_z L$ increases, dropping by 50% for $k_z L = 3 \times 10^{-3}$, and (iii) there is a narrower range of $k_y L$ for unstable modes; for $k_z L = 4.0 \times 10^{-3}$ the waves are unstable only in the region $0.60 < k_y L < 0.93$.

Equation (4) is solved both for the eigenvalues and the eigenfunctions (i.e., the perturbed electrostatic potentials). The eigenfunctions reveal some features of the stabilizing influence of the electron parallel motion. Briefly, the method used to solve (4) is the following: The asymptotic form of the solution is assumed to be WKB type

$$\hat{\phi}_\infty \sim \frac{1}{Q^{1/4}} \exp\left[- \int \sqrt{Q}\, dx\right]$$

where Q is the coefficient of $\hat{\phi}$ in (4). For a given value of $k_y L$, we search for an eigenvalue ω such that integration of (4) with $\hat{\phi}_\infty$ as the boundary condition yields a localized solution whose logarithmic derivative $\partial/\partial x(\ell n\, \hat{\phi})$ is continuous across, for example, the origin. We note that the electron terms are competing with $-k_y^2$ in (4) at $x = \infty$; if the electron terms are dominating Q at $x = \infty$ then Q could become positive, leading to oscillatory type solution at $x = \infty$. On the other hand, if the electron terms are not large such that Q at $x = \infty$ is negative, that is $-k_y^2$, the WKB solutions are exponentially decaying.

Figure 4 shows the solutions of (4) for $k_zL = 3.0 \times 10^{-3}$ and $k_yL = 0.7$ for which $\gamma/(V_0/L) = 0.09$ and $\omega_r/(V_0/L) = 0$. We see that both the real (A) and imaginary (B) parts are exponentially decaying at $x = \pm\infty$ and are fairly localized near the origin where the velocity reversal is taking place. The localization width is proportional to the thickness of the shear layer, $\Delta x \sim 4L$. The value of Q at $x = \infty$ is negative but greater than $-k_y^2L^2$. In Figure 4b we show the full perturbed electrostatic potential, $\hat{\Phi} \sim \hat{\phi}(x) \exp[ik_yy]$ using the wave function, $\hat{\phi}(x)$, for the maximally growing mode of Figure 4a. The dashed (solid) lines indicate the relative negative (positive) values of $\hat{\Phi}$. We note that these values correspond to positive and negative vorticities of the fluid; the phase of $\hat{\Phi}$ is such that it allows mixing of the fluid with different vorticities thereby causing the perturbation to grow.

In Figure 5 we show the eigenfunctions for the lower cutoff wavenumber $k_yL = 0.51$ corresponding to curve D of Fig. 3 for which $k_zL = 3.0 \times 10^{-3}$; the growth rate and the real frequency for this case are $\gamma/(V_0/L) = 0$ and $\omega_r/(V_0/L) = -0.057$, respectively. In this case we find that Q at $x = -\infty$ is positive showing oscillatory type solutions, while Q is negative at $x = \infty$ giving rise to an exponentially decaying solution. In comparison with Figure 4a, this figure shows that electrons are convecting energy away from the localized region in the shear layer, thus, leading to stablization of the instability; this point is discussed in detail in the next section. In Figure 5b we show the corresponding $\hat{\Phi}$. Comparison of Figs. 4b and 5b shows that the parallel motion of electrons changes the phase of Φ and the finite real frequency introduces a relative velocity to the wave preventing any mixing of the fluid in the $x < 0$ and $x > 0$ regions, thus causing the perturbation not to grow.

c. Energy Principle

In this section we present energy principle arguments to show the effects of electron parallel dynamics. The Kelvin-Helmholtz instability arises basically due to the conversion of available kinetic energy of relative motion of the equilibrium flow into wave energy. In this case the electrons are confined to the plane perpendicular to the magnetic field and only $\underset{\sim}{E} \times \underset{\sim}{B}$ drift. If we allow the electrons to move along the magnetic field, the available kinetic energy in the equilibrium flow has to overcome the parallel electron compressional energy for instability to occur.

We can quantify these arguments and obtain a stability condition as follows. We consider the interchange of two neighboring volume elements at x and x + δx which are moving with velocities V_0 and $V_0 + \delta V$, respectively. The available kinetic energy is

$$\delta T = \frac{1}{2} nm_i [V_0^2 + (V_0 + \delta V)^2 - \frac{1}{2}(2V_0 + \delta V)^2] = \frac{1}{4} nm_i (\delta V)^2 \qquad (10)$$

where m_i is the mass of the ions and n is the density. The compressional energy associated with fluid moving a distance δz (i.e., parallel to $\underset{\sim}{B}$) is given by

$$\delta W = < \int_0^{\delta z} nm_e \delta \dot{v}_z dz > = \frac{1}{2} nm_e < \delta \dot{v}_z \delta z > \qquad (11)$$

where m_e is the electron mass and

$$\delta \dot{v}_z = - (e/m_e) \partial \hat{\phi} / \partial z \qquad (12)$$

and < > denotes time averaged quantities. Since the perturbed electron motion in the perpendicular plane is predominantly $\underset{\sim}{E} \times \underset{\sim}{B}$ motion, we have

$$\delta v_x = - (c/B) \partial \hat{\phi} / \partial y \ . \qquad (13)$$

From (12) and (13) we have

$$\frac{\delta v_x}{\delta \dot{v}_z} = \frac{k_y}{k_z \Omega_e} \qquad (14)$$

and

$$\frac{\delta x}{\delta v_z} = \frac{k_y}{k_z \Omega_e} \; . \qquad (15)$$

Substituting (14) in (11) we obtain

$$\delta W = \frac{1}{2}\, nm_e \left(\frac{k_z \Omega_e}{k_y}\right) \langle \delta z\ \delta v_x \rangle = \frac{1}{2}\, nm_e \left(\frac{k_z \Omega_e}{k_y}\right) \langle \delta x\ \delta v_z \rangle$$

from which the compressional energy is given as

$$\delta W = \frac{1}{2}\, nm_e \left(\frac{k_z \Omega_e}{k_y}\right)^2 (\delta x)^2 \; . \qquad (16)$$

From (10) and (16) the stability condition, $\delta T < \delta W$, is

$$\frac{1}{4}\, nm_i (\delta V)^2 < \frac{1}{2}\, nm_e \left(\frac{k_z \Omega_e}{k_y}\right)^2 (\delta x)^2 \qquad (17)$$

or

$$\left(\frac{\partial V_0}{\partial x}\right)^2 < 2\, (\Omega_e \Omega_i)(k_z^2/k_y)^2 \qquad (18)$$

since $\delta V = (\partial V_0/\partial x)\ \delta x$. Thus, the stability condition can be written as

$$\frac{k_z^2}{k_y^2} > \frac{1}{2} \left(\frac{V_0}{L\omega_{\ell h}}\right)^2 \qquad (19)$$

which agrees with (8) derived in Section II.

IV. COLLISIONAL REGIME

In this section we study the effects of electron collisions on the transverse Kelvin-Helmholtz instability. In Fig. 6 we plot the normalized growth rate $(\gamma/(V_0/L))$ as a function of k_zL for $V_d = 0$, $V_0/v_e = 0.01$, $k_zL = 1.0 \times 10^{-3}$, and $k_yL = 0.6$. The variation of the normalized growth rate for $\nu_e = 0$ is shown by curve A, and for $\nu_e = 100\ (V_0/L)$ by curve B. We note that the electron collisions introduce a new scale size, the collisional mean free path

λ_{mfp} ($= v_e/\nu_e$). Since $\lambda_{mfp} << \lambda_z$, electron motion parallel to $\underset{\sim}{B}$ is inhibited, and as a consequence the k_zL domain over which the modes are now unstable is much larger. This is shown clearly in Fig. 6, where the collisional growth rate (curve B) goes to zero for $k_zL > 0.07$, while the collisionless growth rate (curve A) drops to zero sharply for $k_zL > 0.005$.

In addition, curve B of Fig. 2 shows the collosional growth rate versus the coupling parameter α. For the range of α considered, electron collisions make the growth rate almost independent of α. The electron collisions thus have a destabilizing influence on Kelvin-Helmholtz instability. This would in turn alter the stability boundary in such a way that the basic Kelvin-Helmholtz mode is unstable in a larger k_z domain than the collisionless case.

V. SUMMARY AND DISCUSSION

In this paper we have considered the effect of parallel electron motion on the transverse Kelvin-Helmholtz instability as a first step toward a three-dimensional analysis of the nature of convective flows in the high latitude ionosphere. In the collisionless regime ($\nu_e = 0$), we show that for an inhomogeneous velocity profile given by $V(x) = V_0 \tanh(x/L)$ the Kelvin-Helmholtz instability is stable for $k_z/k_y > (V_0/L\omega_{\ell h})k_yL(1 - k_y^2L^2)^{1/2}$ in the limit $\omega - k_yV_0 >> k_zv_e$. We further show that the physical mechanism for stabilization is the compressional energy given to the electrons parallel to $\underset{\sim}{B}$. In the collisional regime ($\nu_e \neq 0$), we show that the Kelvin-Helmholtz instability is not as easily stabilized by finite k_z effects; this is attributed to the inhibition of parallel electron motion for $\lambda_{mfp} << \lambda_z$.

We apply this stability criterion for parameters relevant to the high latitude ionosphere. We take $\omega_{\ell h} = 3.2 \times 10^4$ sec^{-1}, $k_yL = 0.5$, and assume $k_z/k_y = (\sigma_\perp/\sigma_\parallel)^{1/2}$ where $\sigma_\perp$ and $\sigma_\parallel$ are the perpendicular and parallel conductivities, respectively (Farley, 1959). The value of $\sigma_\perp/\sigma_\parallel$ varies with altitude in the ionosphere and is generally taken to be in the range 10^{-5} (lower F region) - 10^{-9} (topside F region). Thus, for instability to occur it is required that $V_0/L > 2.3$ sec^{-1} in the topside region. Velocity shears have been reported in the range $V_0/L = 0.5 - 20$ Hz (Kelley and Carlson, 1977; Earle and Kelley, 1986) so that the instability could locally be active in the topside F region. In the lower F region, where $\nu_e \sim 10^2 - 10^3$ sec^{-1} and $V_0/v_e =$

10^{-2} velocity shears of order 1 - 10 sec^{-1} are needed to drive the Kelvin-Helmholtz modes unstable (e.g., for $k_y L = 0.6$ and $k_z L = 0.01$). However, these conclusions are predicated on the important assumption that $k_z/k_y \approx (\sigma_\perp/\sigma_\parallel)^{1/2}$ which may not be the case. Recent analyses of unstable flute modes in barium clouds (i.e., the gradient drift instability) (Sperling et al., 1984; Drake et al., 1985) indicate that $k_z/k_y \approx 0$ within the unstable cloud. A more detailed three-dimensional analysis (Drake and Huba, 1987) is needed to assess the role of the Kelvin-Helmholtz instability in high latitude dynamics. Finally, we add that we have neglected ion-neutral collisions in the analysis; these can be important in the lower F region. In a forthcoming study (Mitchell et al., 1987) we will present analytical and computational results demonstrating the effect of ion-neutral collisions on the transverse Kelvin-Helmholtz instability. In the linear regime, it is found that ion-neutral collisions have a stabilizing influence on the instability.

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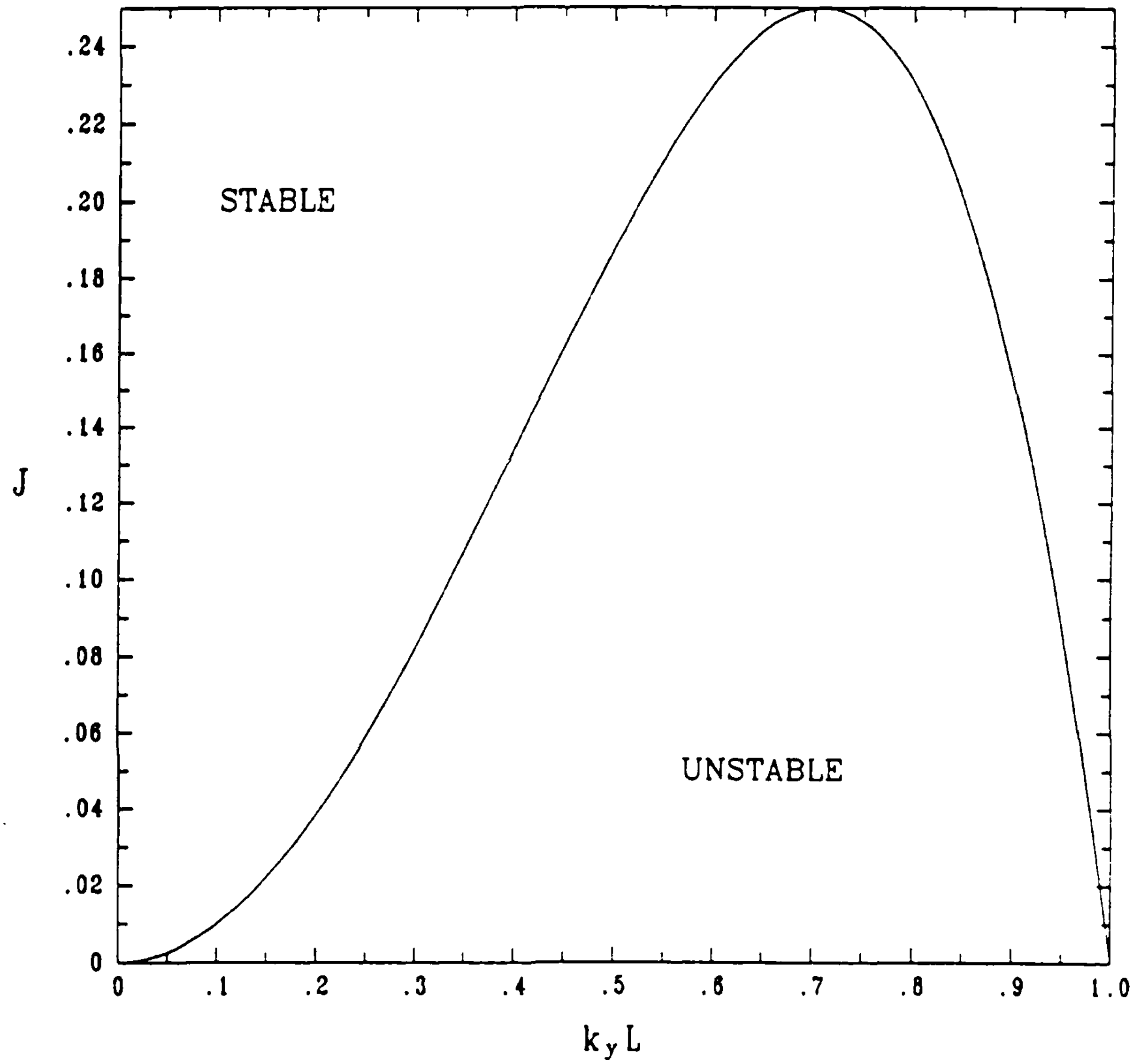

Fig. 1) Stability boundary in the fluid limit. The curve shows the marginal stability boundary ($\gamma = 0$) of $J = \left(L^2 \omega_{\ell h}^2 / V_0^2\right)\left(k_z^2/k_y^2\right)$ vs $k_y L$.

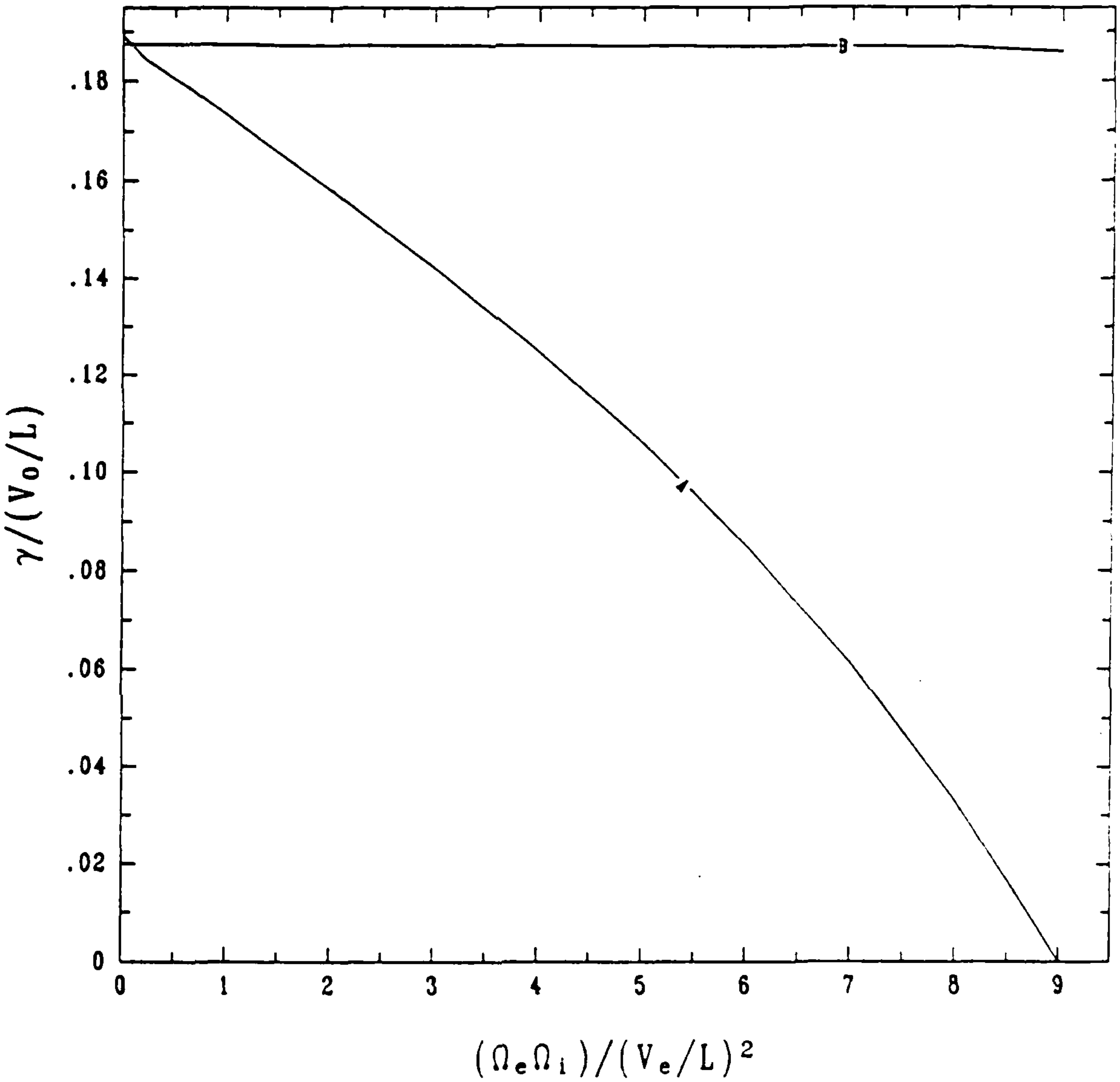

Fig. 2) Plot of the normalized growth rate $\gamma/(V_0/L)$ vs α $(= L^2\omega_{\ell h}^2/v_e^2)$. The parameters used are $k_yL = 0.5$, $k_z/k_y = 1.0 \times 10^{-3}$, $V_0/v_e = 0.01$, and $V_d = 0$. Curve A is for the collisionles case ($\nu_e = 0$), while curve B is for the collisional case ($\nu_e = 100\ V_0/L$).

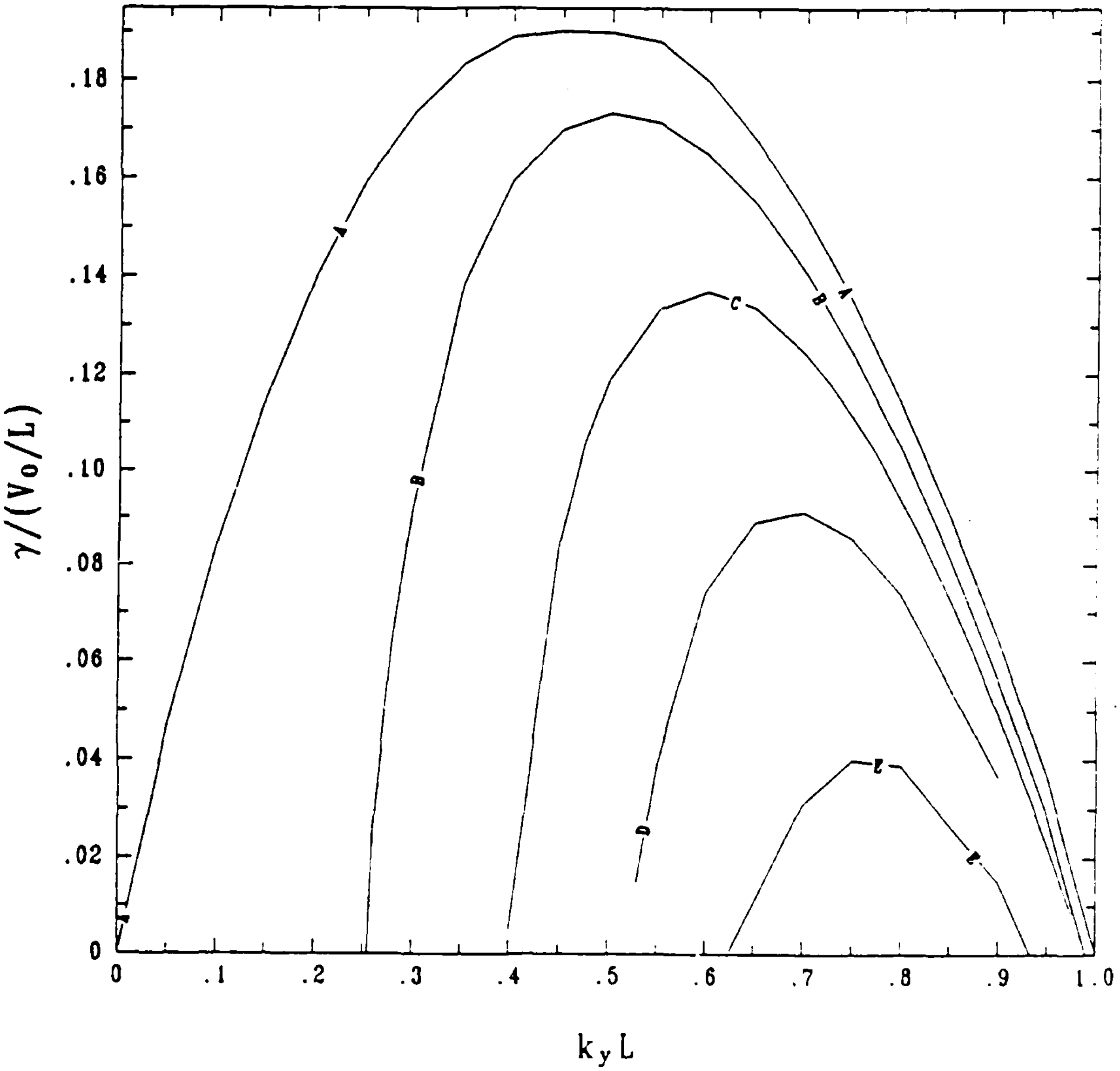

Fig. 3) Plot of the normalized growth rate $\gamma/(V_0/L)$ vs k_yL for V_0/v_e = 0.01, α = 1, V_d = 0 and k_zL = 0, 1 x 10^{-3}, 2 x 10^{-3}, and 4 x 10^{-3}, labelled A - E, respectively.

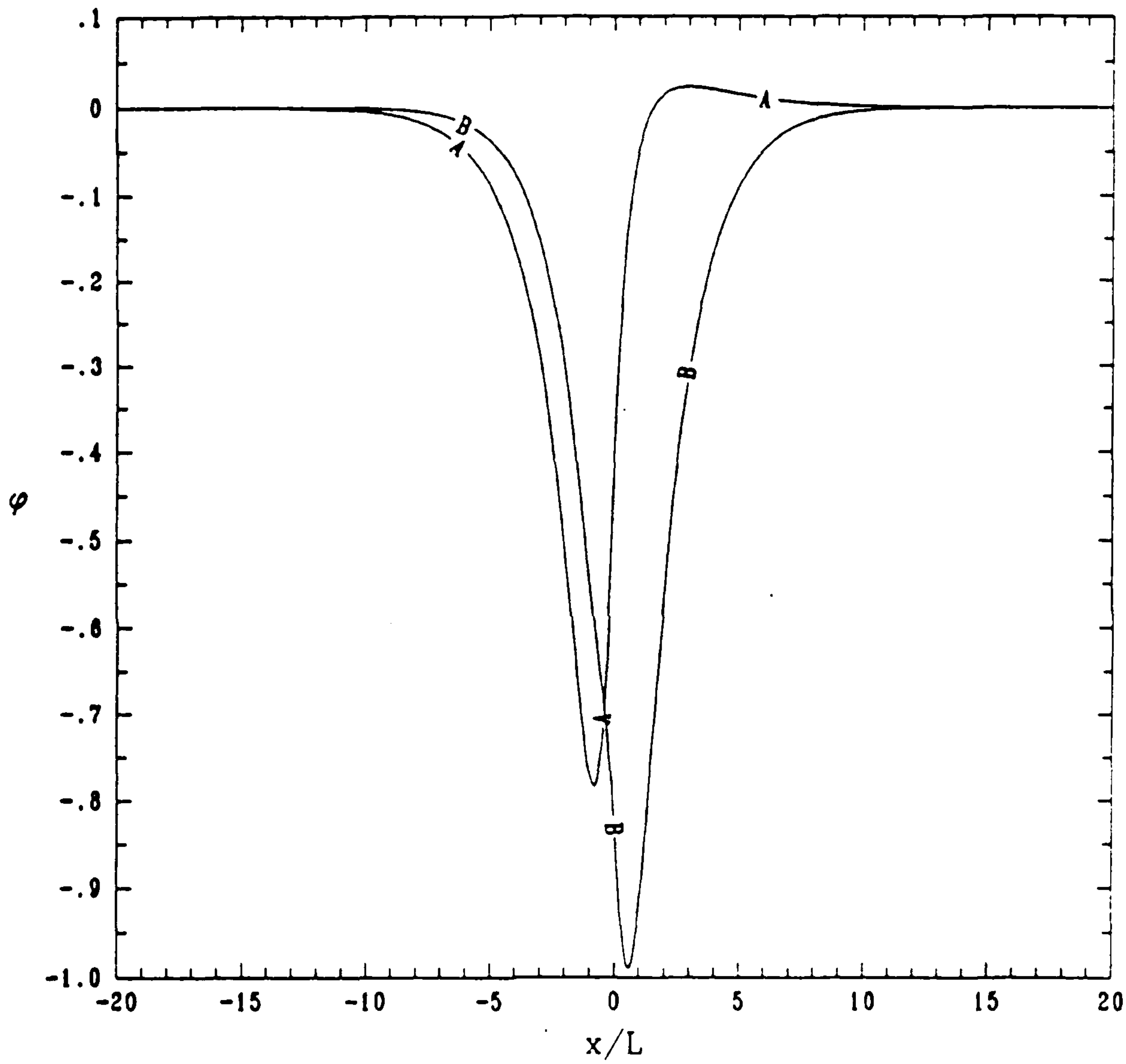

Fig. 4) Eigenfunction for the mode corresponding to curve D of Fig. 3 for $k_yL = 0.7$ ($\gamma = 0.09$). (a) Curves A and B refer to the real and imaginary parts of $\hat{\phi}$, respectively. (b) Contour plots of $\hat{\Phi} \sim \hat{\phi}(x) \exp[ik_y y]$.

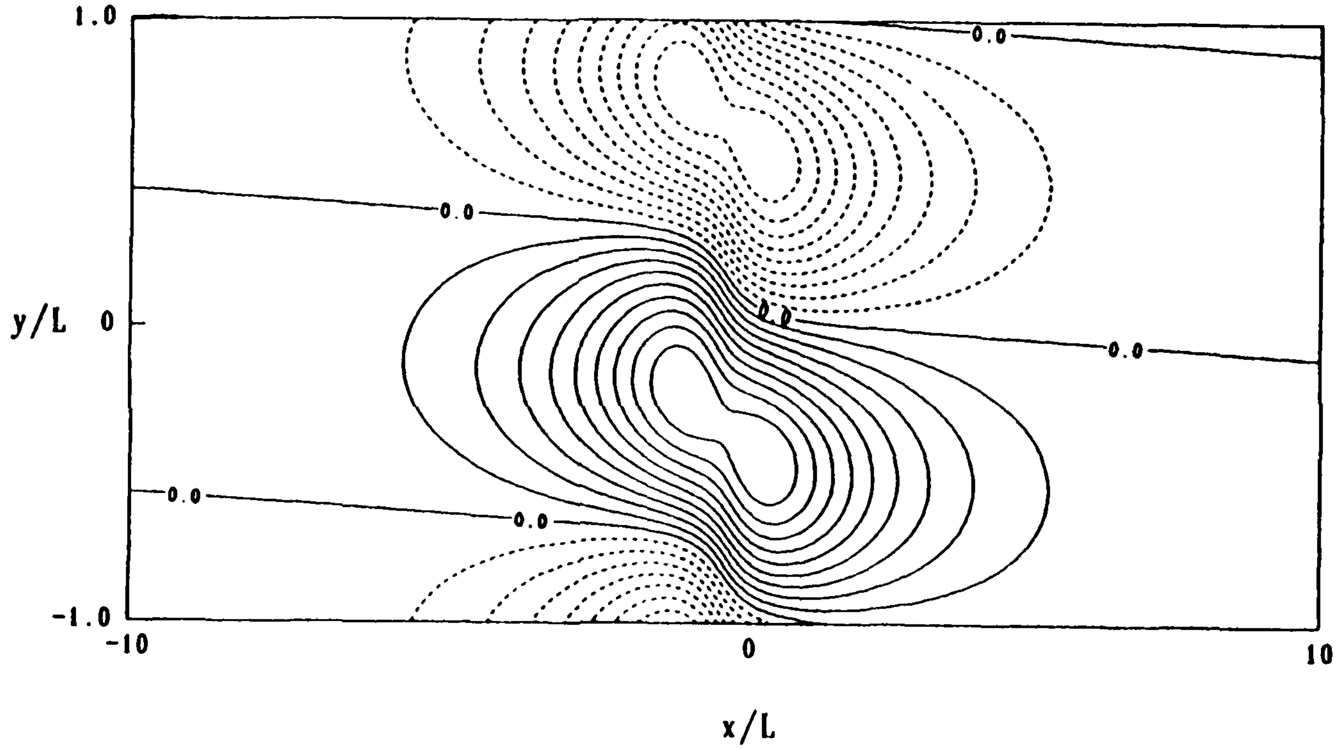

Fig. 4) (Cont) Eigenfunction for the mode corresponding to curve D of Fig. 3 for $k_y L = 0.7$ ($\gamma = 0.09$). (a) Curves A and B refer to the real and imaginary parts of $\hat{\phi}$, respectively. (b) Contour plot of $\tilde{\phi} \sim \hat{\phi}(x) \exp[ik_y y]$.

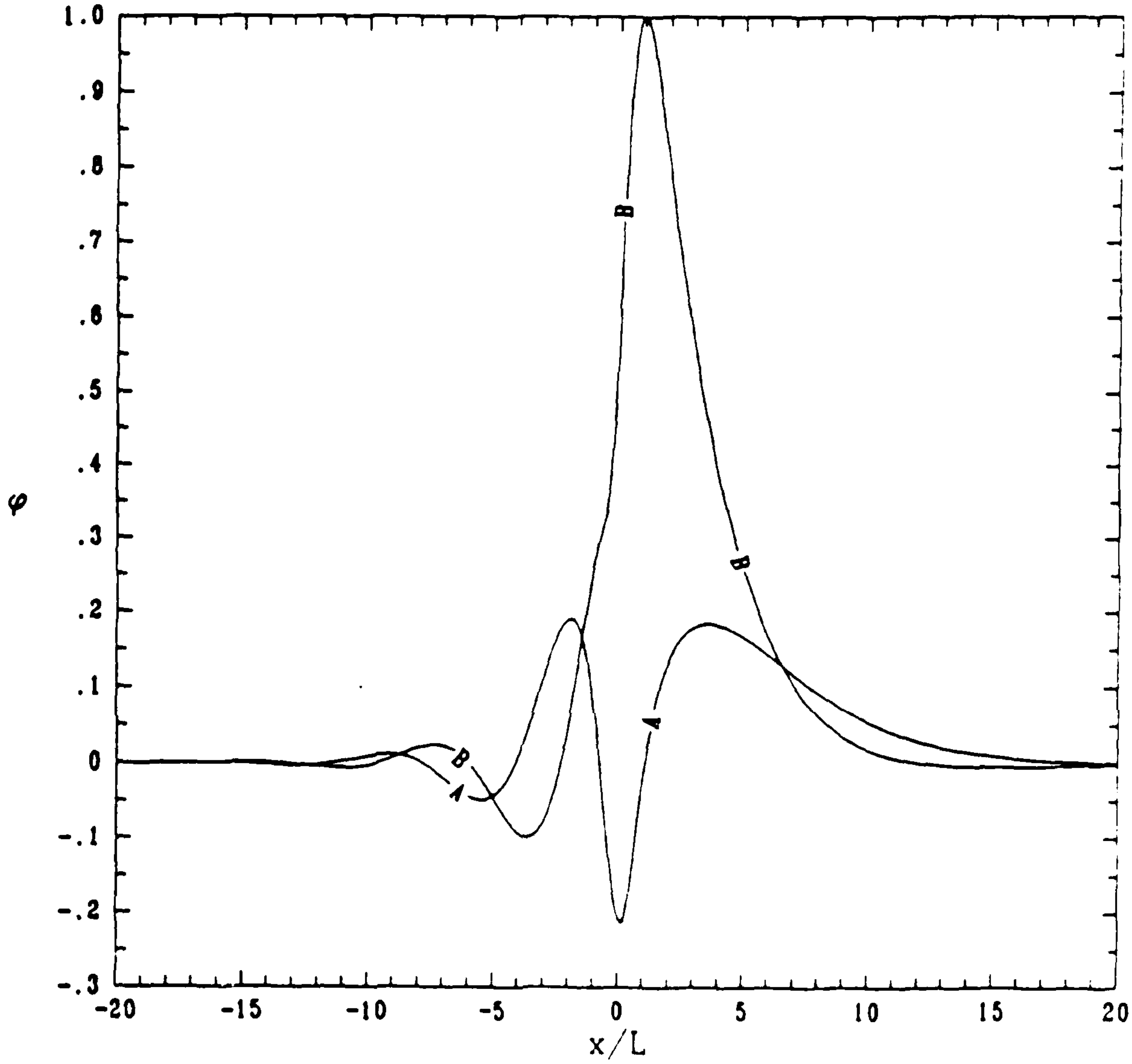

Fig. 5) Eigenfunction for the mode corresponding to Curve D of Fig. 3 for $k_yL = 0.48$ ($\gamma = 0$). (a) Curves A and B refer to the real and imaginary parts of $\hat{\phi}$, respectively. (b) Contour plots of $\hat{\Phi} = \hat{\phi}(x) \exp[ik_y y]$.

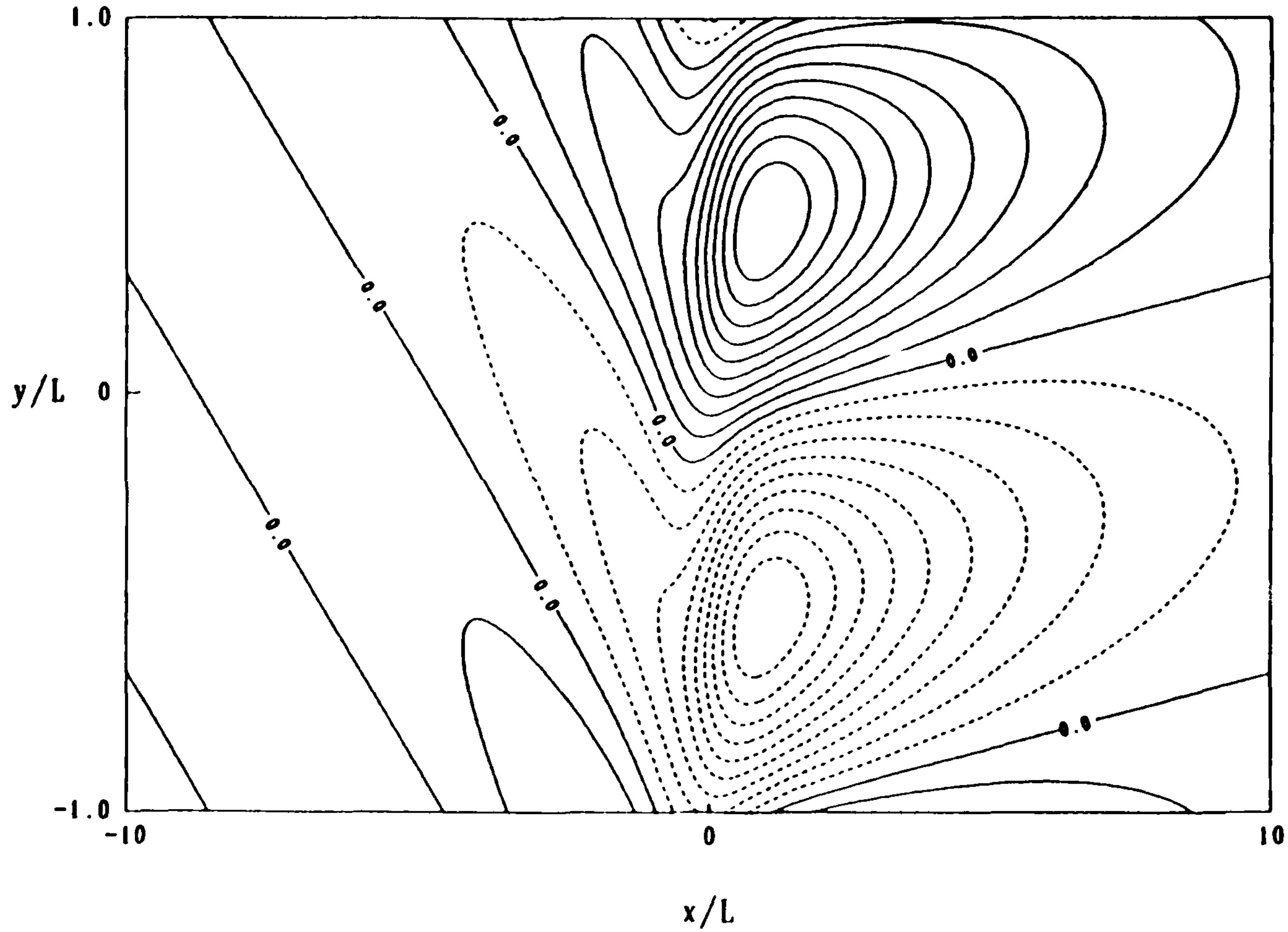

Fig. 5) (Cont) Eigenfunction for the mode corresponding to Curve D of Fig. 3 for $k_y L = 0.48$ ($\gamma = 0$). (a) Curves A and B refer to the real and imaginary parts of $\hat{\phi}$, respectively. (b) Contour plots of $\tilde{\phi} = \hat{\phi}(x) \exp[ik_y y]$.

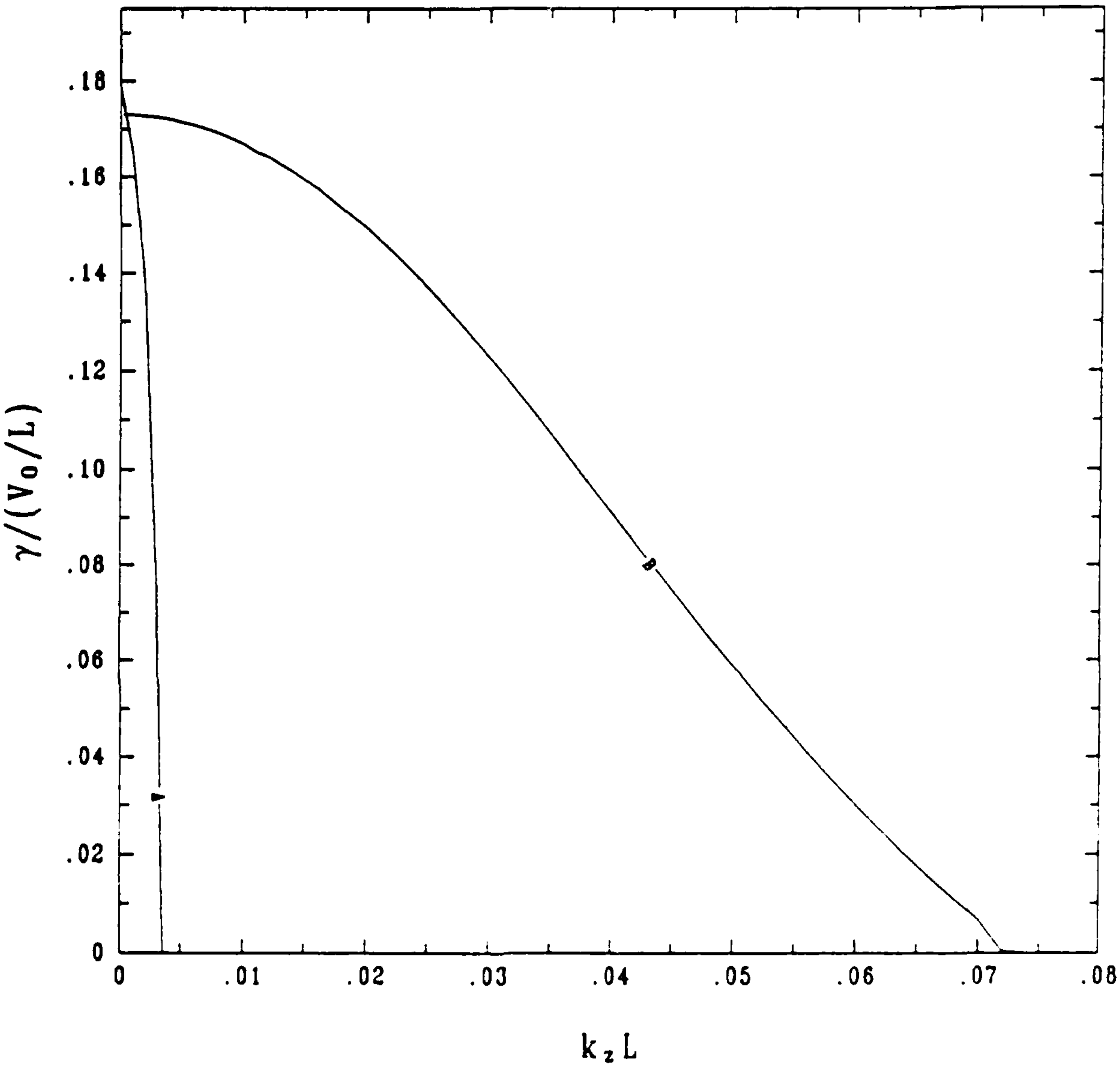

Fig. 6) Plot of the normalized growth rate $\gamma/(V_0/L)$ vs $k_z L$ for V_d = 0, $k_y L$ = 0.5, V_0/V_e = 0.01, and α = 1. Curve A corresponds to the case ν_e = 0 and curve B to ν_e = 100 (V_0/L), respectively.

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